T violation in radiative β decay and electric dipole moments

W. G. Dekens and K. K. Vos

Van Swinderen Institute, Faculty of Mathematics and Natural Sciences,

University of Groningen, Nijenborgh 4,

NL-9747 AG Groningen, The Netherlands

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Abstract

In radiative β decay, T violation can be studied through a spin-independent T-odd correlation. We consider contributions to this correlation by beyond the standard model (BSM) sources of T-violation, arising above the electroweak scale. At the same time such sources, parametrized by dimension-6 operators, can induce electric dipole moments (EDMs). As a consequence, the manifestations of the T-odd BSM physics in radiative β decay and EDMs are not independent. Here we exploit this connection to show that current EDM bounds already strongly constrain the spin-independent T-odd correlation in radiative β decay.

I. INTRODUCTION

The Standard Model of particle physics (SM) cannot account for the baryon asymmetry of the universe [1–3], and additional sources of CP violation might be expected to arise beyond the SM (BSM). Searches for additional time-reversal (T) violation, and equivalently CP violation, are, therefore, promising probes of BSM physics. Especially interesting are observables with a very low SM background, such as the electric dipole moments (EDMs) of hadrons, nuclei, atoms, and molecules.

In β decay, T violation is probed by the triple-correlation coefficients D and R [4]. However, these observables are not independent from EDM measurements. In fact, the stringent neutron EDM limit bounds D more than an order of magnitude better than current β-decay experiments [5–8]. Molecular and atomic EDMs constrain scalar and tensor electron-nucleon couplings, which leads to strong constraints on the R coefficient. These constraints are several order of magnitude better than the current best β-decay bounds [9–12].

In radiative decays it is possible to study spin-independent T-odd triple-correlations [13–15], which are not present in β decay. In this paper, we consider such a correlation in radiative β decay generated by high-energy BSM sources of CP violation. As in β decay, we find that this T-odd correlation and EDMs are connected, which allows EDM bounds to strongly constrain the spin-independent T-odd correlation.

We work in an effective field theory (EFT) framework in which dimension-6 operators parametrize the new sources of CP violation. We first discuss these operators. Then we consider their contributions to radiative β decay in Section III, while discussing the contribution of these operators to the EDM in section IV. Finally, we give the current EDM bounds on these operators and end with a brief discussion.

II. FORMALISM

We consider the effects of new T-violating physics on the correlation $K\vec{p}_\nu \cdot (\vec{p}_e \times \vec{k})$, where $\vec{k}$ is the photon momentum, and neglect the small T-violating SM contributions generated by the CP-odd phase of the CKM matrix and the QCD θ-term [16]. Besides these true SM T-odd sources, there are also electromagnetic final-state interactions (FSI) that mimic T-violation and that also contribute to the triple-correlation (similar FSI contribute to D and

R). These FSI have been studied for the neutron, ^{19}Ne, and ^{35}Ar [14, 15, 17], and contribute to the T-odd asymmetry at $\mathcal{O}(10^{-3}) - \mathcal{O}(10^{-5})$, depending on the detectable photon energy and the used isotope.

The effects of new T-violating physics, arising at a high scale Λ, can be studied in an EFT framework. At low energies, the new physics is effectively described by higher-dimensional operators. We consider dimension-six operators, for which the complete set of gauge-invariant operators has been derived in Ref. [18]. We divide the operators relevant for radiative β decay into two groups.

(i) The first group consists of four-fermion operators that also contribute to β decay [18, 19]. The relevant part of the effective β-decay Lagrangian is [20]

$$\mathcal{L}_{S,P,T}^{(\text{eff})} = \frac{-4G_F}{\sqrt{2}} \sum_{\epsilon,\delta=L,R} \left\{ A_{\epsilon\delta}\, \bar{e}\nu_e^\epsilon \cdot \bar{u}d_\delta + \alpha_\epsilon\, \bar{e}\frac{\sigma^{\mu\nu}}{\sqrt{2}}\nu_e^\epsilon \cdot \bar{u}\frac{\sigma_{\mu\nu}}{\sqrt{2}}d_\epsilon \right\} + \text{h.c.}, \tag{1}$$

where we have set $V_{ud} = 1$ for convenience. G_F is the Fermi coupling constant and we sum over the chirality (L, R) of the final states. These four-fermion operators modify the $V - A$ coupling of the SM, by generating scalar/pseudoscalar (A) and tensor (α) couplings [4, 21]. Besides contributing to β decay, the operators in Eq. (1) also contribute to radiative β decay after being dressed with bremsstrahlung photons.

(ii) The second group of T-violating operators is given in Table I[1]. At the scale of new physics, Λ, the relevant terms for radiative β decay, are

$$\mathcal{L}_6 = C_{\varphi\tilde{W}B}(\Lambda)\frac{gc_w v^2}{2}i\varepsilon^{\mu\nu\alpha\beta}W_\mu^+ W_\nu^- F_{\alpha\beta} + C_{\varphi ud}(\Lambda)\frac{v^2 g}{2\sqrt{2}}\bar{u}_R\gamma^\mu d_R W_\mu^+$$

$$+ 2vC_{uW}(\Lambda)\,(\bar{d}_L\sigma^{\mu\nu}\overleftrightarrow{D}_\nu u_R)W_\mu^- + 2vC_{dW}(\Lambda)\,(\bar{u}_L\sigma^{\mu\nu}\overleftrightarrow{D}_\nu d_R)W_\mu^+$$

$$+ \text{h.c.} + \ldots, \tag{2}$$

where $v \approx 246\,\text{GeV}$ is the vacuum expectation value of the Higgs field $\langle\varphi\rangle = \frac{1}{\sqrt{2}}v$, the photon field is denoted by A_μ and $s_w = \sin\theta_w$ is the sine of the Weinberg angle ($c_w = \cos\theta_w$). The covariant derivative $D_\mu = \partial_\mu - is_w g q_f A_\nu$, where q_f is the charge of the fermion. C_X is the coupling constant associated with the operator Q_X defined in Table I.

Fig. 1 shows how these operators contribute to radiative β decay. At low energies, $\mu \approx$

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[1] In principle, the operator $Q_{eW} = (\bar{l}\sigma^{\mu\nu}e)\tau^I\varphi W_{\mu\nu}^I$ also contributes to radiative β decay, however, it does not contribute to K at leading recoil order.

$Q_{\varphi ud}$	$i(\tilde{\varphi}^\dagger D_\mu \varphi)(\bar{u}\gamma^\mu d)$
$Q_{\varphi \tilde{W} B}$	$\varphi^\dagger \tau^I \varphi \tilde{W}^I_{\mu\nu} B^{\mu\nu}$
Q_{uW}	$(\bar{q}\sigma^{\mu\nu}\tau^I \tilde{\varphi}\, u)W^I_{\mu\nu}$
Q_{dW}	$(\bar{q}\sigma^{\mu\nu}\tau^I \varphi\, d)W^I_{\mu\nu}$

TABLE I. Dimension-six operators that contribute to T-violating radiative β decay. Here τ^I are the Pauli matrices, φ is the Higgs doublet and $\tilde{\varphi} = i\tau_2 \varphi^*$. Furthermore, $D_\mu = \partial_\mu - i\frac{g}{2}\tau^I W^I_\mu - i\frac{g'}{2}B_\mu$ is the covariant derivative of the Higgs doublet, while $W^I_{\mu\nu} = \partial_\mu W^I_\nu - \partial_\nu W^I_\mu + g\varepsilon^{IJK}W^J_\mu W^K_\nu$ and $B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$ are the field strengths of the $SU(2)$ and $U(1)_Y$ gauge fields respectively. Finally, the duals of the field strengths are $\tilde{X}_{\mu\nu} = \varepsilon_{\mu\nu\alpha\beta}X^{\alpha\beta}$, where $\varepsilon^{0123} = +1$.

$1\,\text{GeV}$, after integrating out the $W^\pm$ boson, we obtain

$$
\begin{aligned}
\mathcal{L}_6^{\text{eff}} = & -\frac{8ic_w}{gv^2}V_{ud}\,\text{Re}\,C_{\varphi \tilde{W} B}(\Lambda)\,\varepsilon^{\mu\nu\alpha\beta}(\bar{u}_L\gamma_\mu d_L)(\bar{e}_L\gamma_\nu \nu_L)F_{\alpha\beta} + \frac{1}{M_W^2}C_{\varphi ud}(\Lambda)(\bar{u}_R\gamma_\mu d_R)\Gamma^{\mu\nu}(\bar{e}_L\gamma_\nu \nu_L) \\
& -\frac{8is_w}{\sqrt{2}v}\eta_{qW}C^*_{Wu}(\Lambda)\,(\bar{u}_R\sigma^{\mu\nu}d_L)(\bar{e}_L\gamma_\mu \nu_L)A_\nu - \frac{8is_w}{\sqrt{2}v}\eta_{qW}C_{Wd}(\Lambda)\,(\bar{u}_L\sigma^{\mu\nu}d_R)(\bar{e}_L\gamma_\mu \nu_L)A_\nu \\
& +\text{h.c.} + \dots,
\end{aligned}
\tag{3}
$$

where $\Gamma^{\mu\nu} = g^{\mu\nu}D^2 - D^\nu D^\mu - igs_w F^{\mu\nu}$, whose leading contribution to K arises from Fig. 1a. Furthermore, $\eta_{qW} = \left(\frac{\alpha_s(\Lambda)}{\alpha_s(m_t)}\right)^{4/21}\left(\frac{\alpha_s(m_t)}{\alpha_s(m_b)}\right)^{4/23}\left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)}\right)^{4/25}\left(\frac{\alpha_s(m_c)}{\alpha_s(\mu)}\right)^{4/27}$ is a running factor (numerically, $\eta_{qW} = 0.39\,(0.33)$ for $\Lambda = 1\,(10)\,\text{TeV}$), arising from the QCD renormalization of the Q_{qW} operators [22, 23]. The dots represent terms which are necessary to maintain gauge invariance, but that do not contribute to K at leading recoil order.

The first term in Eq. (3) is similar to the interaction studied in Eq. (2) of Ref. [15]. Although we find that such a term is not T-violating when it arises from a pseudo-Chern-Simons term (*i.e.* Eq. (1) in Ref. [15]), it is clear that it can be generated by BSM physics such as $Q_{\varphi \tilde{W} B}$.

III. T-VIOLATING RADIATIVE β DECAY

The new sources of T violation contribute to the radiative β decay rate

$$
d\Gamma = 32e^2 G_F^2 M_n M_p d\Gamma_0 \left[K\vec{p}_\nu \cdot (\vec{p}_e \times \vec{k}) + \cdots \right] ,
\tag{4}
$$

where $M_{n,p}$ are the neutron and proton masses and $d\Gamma_0$ contains the integral over the phase space. The dots represent higher-order recoil terms as well as T-even terms that are present
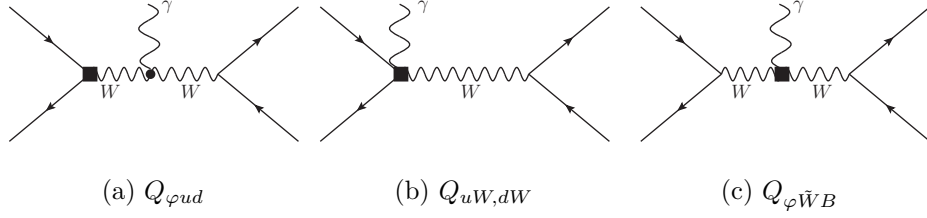
4

(a) $Q_{\varphi ud}$ (b) $Q_{uW,dW}$ (c) $Q_{\varphi \tilde{W}B}$

FIG. 1. The relevant T-violating effective interactions for radiative β decay. The box indicates one of the BSM interactions of Table I, the dot indicates the SM coupling of two W bosons and a photon.

in the SM [24]. The K coefficient can be inferred from the asymmetry [14, 15]

$$\mathcal{A} = \frac{\Gamma^+ - \Gamma^-}{\Gamma^+ + \Gamma^-} \, , \tag{5}$$

where Γ^+ corresponds to the ϕ_ν range $[0, \pi]$ and Γ^- to the ϕ_ν range $[\pi, 2\pi]$ if $\vec{p}_e$ is in the $\hat{z}$ direction, such that $\vec{k}$ and $\vec{p}_e$ fix the $\hat{z} - \hat{x}$ plane [14, 15]. The asymmetry depends on the Q-value of the interaction and on the threshold energy of the photon detector $\omega_{\min}$, typically in the range of MeV. The asymmetry grows with increasing $\omega_{\min}$, following Ref. [15] we evaluate the integrals at $\omega_{\min} = 0.3$MeV. We discuss the form of K and the contribution to $\mathcal{A}$ for the two groups of operators.

(i) The first group contributes to radiative β decay after being dressed with bremsstrahlung photons [15]. For neutron decay,

$$K = 2\frac{1}{M_p}\frac{1}{k \cdot p_e}\text{Im}\left[g_T\alpha_L(g_S^*A_L^* + g_P^*A_L'^*) - g_T\alpha_R(g_S^*A_R^* + g_P^*A_R'^*)\right] \, , \tag{6}$$

where $A_L \equiv A_{LL} + A_{LR}$, $A_R \equiv A_{RR} + A_{RL}$, $A_L' \equiv A_{LL} - A_{LR}$ and $A_R' \equiv A_{RR} - A_{RL}$. The couplings g_Λ are defined by $\langle p|\bar{u}\Gamma d|n\rangle = g_\Gamma \bar{p}\Gamma n$, with $\Gamma = 1, \gamma_5, \gamma_\mu, \gamma_\mu\gamma_5, \sigma^{\mu\nu}$. For $\omega_{\min} = 0.3$ MeV, the asymmetry is

$$\mathcal{A} = 2.1 \times 10^{-5}\text{Im}\left[g_T\alpha_L(g_S^*A_L^* + g_P^*A_L'^*) - g_T\alpha_R(g_S^*A_R^* + g_P^*A_R'^*)\right] \, , \tag{7}$$

which is in part small due to the nucleon mass suppression in Eq. (6). The asymmetry only contains quadratic couplings, which also appear in the R correlation. The constraints on these couplings from EDM measurements and the R coefficient range from $\mathcal{O}(10^{-3})-\mathcal{O}(10^{-6})$ [11, 12]. Improving these bounds in radiative β decay would require a measurement of the asymmetry to better than 10^{-11}.
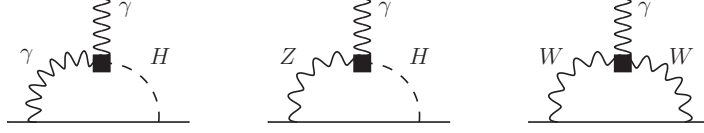
FIG. 2. One-loop contributions of $Q_{\varphi \tilde{W} B}$ to the quark and electron EDMs.

(ii) At leading order, the interactions in Eq. (3) give

$$
K = -16 \frac{c_w}{eg} \frac{E_e}{k \cdot p_e} (g_A^2 + g_V^2) \mathrm{Re}\, C_{\varphi \tilde{W} B} - 8 s_w \frac{1}{k \cdot p_e} \frac{\sqrt{2} M_W}{eg} g_A g_T \eta_{qW} \mathrm{Im}(C_{Wd}^* + C_{Wu})
$$
$$
+ \frac{s_w}{eg} \mathrm{Im} C_{\varphi ud} \left(8 \frac{E_e}{k \cdot p_e} (g_A^2 - g_V^2) + 4 \frac{1}{\omega} (g_A^2 + g_V^2) \right) . \tag{8}
$$

For $\omega_{\min} = 0.3$ MeV, the asymmetry for neutron radiative β decay is

$$
\mathcal{A} = -2 \times 10^{-11}\, \mathrm{Re}\, C_{\varphi \tilde{W} B}' + 2 \times 10^{-7}\, \mathrm{Im}(C_{Wd}'^* + C_{Wu}')
$$
$$
+ 4 \times 10^{-12}\, \mathrm{Im} C_{\varphi ud}' , \tag{9}
$$

in terms of the couplings at $\Lambda = 1\,\mathrm{TeV}$. We used $g_A = 1.27$, $s_w^2 = 0.23$, $g = 0.64$, $g_T \sim 1$ and $M_W = 80.4$ GeV. For clarity we have redefined $C' \equiv v^2 C$, such that the couplings C' are dimensionless. Clearly, the contribution of these operators to the asymmetry is rather small. However, the sensitivity of the asymmetry to the T-odd BSM sources can be improved by choosing isotopes with larger Q-values. For ^{37}K [25], we find, for example, that it is 20 times more sensitive than neutron decay. In the next Section we discuss the stringent limit from EDMs on these couplings.

IV. CONSTRAINTS FROM EDMS

The second class of operators also contribute to the neutron EDM (nEDM) and electron (eEDM). At the scale M_W, these operators induce

$$
\mathcal{L}_{\mathrm{EDM}} = -\frac{i}{2} \sum_{f=u,d,e} d_f e Q_f m_f\, \bar{\psi}_f \sigma^{\mu\nu} \gamma_5 \psi_f F_{\mu\nu} - i\, \mathrm{Im}\, \Xi_1 \left[\bar{u}_R \gamma^\mu d_R\, \bar{d}_L \gamma_\mu u_L - \bar{d}_R \gamma^\mu u_R\, \bar{u}_L \gamma_\mu d_L \right]
$$
$$
- i\, \mathrm{Im}\, \Xi_8 \left[\bar{u}_R \gamma^\mu t^a d_R\, \bar{d}_L \gamma_\mu t^a u_L - \bar{d}_R \gamma^\mu t^a u_R\, \bar{u}_L \gamma_\mu t^a d_L \right], \tag{10}
$$

where $d_{u,d}$ represent the up- and down-quark EDM, d_e is the electron EDM ($d_e^{\mathrm{exp}} \equiv e m_e d_e$), and $\Xi_{1,8}$ are CP-odd four quark operators. The contributions from $C_{\varphi ud}$, $C_{\varphi \tilde{W} B}$, and C_{qW} to the couplings in Eq. (10) are listed in Table II.

	$C_{\varphi ud}(\Lambda)$	$\frac{1}{gg'}\operatorname{Re} C_{\varphi\tilde{W}B}(\Lambda = 1\,(10)\,\text{TeV})$	$\frac{\sqrt{2}vs_w}{em_qQ_q}\operatorname{Im} C_{qW}(\Lambda = 1\,(10)\,\text{TeV})$
$d_u(M_W)$	–	$-6.2\,(-11)\times 10^{-2}$	$-0.80\,(-0.69)$
$d_d(M_W)$	–	$-11\,(-19)\times 10^{-2}$	$0.80\,(0.69)$
$d_e(M_W)$	–	$-5.3\,(-10)\times 10^{-2}$	–
$\Xi_1(M_W)$	1	–	–
$\Xi_8(M_W)$	–	–	–

TABLE II. The couplings in Eq. (10) at the scale M_W in terms of the BSM couplings $C_{\varphi ud}$, $C_{\varphi\tilde{W}B}$, and C_{qW} at the scale $\Lambda = 1\,(10)\,\text{TeV}$. The row headings are given by the column headings multiplied by the corresponding table entry, a dash indicates there is no contribution at leading order in perturbation theory. For more details, see Ref. [27].

$Q_{\varphi ud}$

Table II shows that, at leading order $C_{\varphi ud}$, only contributes to Ξ_1. This interaction is generated after integrating out the $W^{\pm}$ boson through a tree-level diagram. The relevant interaction is similar to Fig. 1b without the photon, and where the $W^{\pm}$ should now be coupled to quarks instead of leptons [5]. As $Q_{\varphi ud}$ does not evolve under QCD renormalization, the relation in Table II is independent of Λ.

$Q_{\varphi\tilde{W}B}$

In contrast, $Q_{\varphi\tilde{W}B}$ does not contribute to the interactions in Eq. (10) at the tree-level, but it induces quark EDMs at the one-loop-level. This operator also contains Higgs $-\gamma\gamma$ and Higgs $-Z\gamma$ interactions. These interactions and the $WW\gamma$ interaction in Eq. (2) contribute to the quark EDMs through the diagrams in Fig. 2 [26, 27]. As a consequence, the operator $Q_{\varphi\tilde{W}B}$ mixes with the quark EDMs when it is evolved from the scale of new physics, Λ, down to M_W. The electron EDM (eEDM) is induced through the same mechanism. The results in Table II take into account both the mixing between $C_{\varphi\tilde{W}B}$ and d_q, and the running of d_q^2 [22, 23].

Q_{uW} and Q_{dW}

Finally, the Q_{qW} operators contribute to $d_{u,d}$ directly, and we have,

$$d_u(\Lambda) = -\frac{\sqrt{2}vs_w}{eQ_um_u}\operatorname{Im} C_{uW}(\Lambda), \qquad d_d(\Lambda) = \frac{\sqrt{2}vs_w}{eQ_dm_d}\operatorname{Im} C_{dW}(\Lambda). \tag{11}$$

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[2] Like $Q_{\varphi ud}$, the operator $Q_{\varphi\tilde{W}B}$ does not evolve under one-loop QCD renormalization.

nEDM	$\mathrm{Im}\, C'_{\varphi ud}(\Lambda)$	$\mathrm{Re}\, C'_{\varphi\tilde{W}B}(\Lambda)$	$\mathrm{Im}\, C'_{qW}(\Lambda)$		eEDM	$\mathrm{Re}\, C'_{\varphi\tilde{W}B}(\Lambda)$
$\Lambda = 1\,\mathrm{TeV}$	1.0×10^{-5}	1.8×10^{-4}	2.1×10^{-10}		$\Lambda = 1\,\mathrm{TeV}$	2.3×10^{-6}
$\Lambda = 10\,\mathrm{TeV}$	1.0×10^{-5}	1.0×10^{-4}	2.2×10^{-10}		$\Lambda = 10\,\mathrm{TeV}$	1.2×10^{-6}

TABLE III. Bounds on the couplings of the dimension-six operators in Table I due to the limits on the neutron and electron EDM. Again $C' \equiv v^2 C$ to make the primed couplings dimensionless. The bounds are shown for two values of the scale of new physics, $\Lambda = 1, 10\,\mathrm{TeV}$. In the cases that an operator contributes in several ways to d_n we present the strongest bound, not taking into account possible cancellations. Only the $Q_{\varphi\tilde{W}B}$ gives rise to a significant eEDM.

After taking into account the running of the quark EDMs, we obtain the results in Table II.

The induced interactions at the scale M_W have to be evolved to the low energies where EDM experiments take place. The renormalization group equations (RGEs) for the quark EDMs and the four-quark operators give [22, 23, 27–29]

$$d_q(M_{\mathrm{QCD}}) = 0.48\, d_q(M_W),$$

$$\mathrm{Im}\,\Xi_1(M_{\mathrm{QCD}}) = 1.1\, \Xi(M_W), \qquad \mathrm{Im}\,\Xi_8(M_{\mathrm{QCD}}) = 1.4\, \Xi(M_W), \tag{12}$$

where $M_{\mathrm{QCD}} \approx 1\,\mathrm{GeV}$ is the QCD scale, while the eEDM does not evolve under one-loop QCD renormalization. For calculation of the nEDM in terms of d_q and $\Xi_{1,8}$ we use the following naive dimensional analysis [30, 31] (NDA) estimates,

$$d_n^{d_q} = \mathcal{O}(eQ_q m_q)\, d_{u,d}(M_{\mathrm{QCD}}), \qquad d_n^{\Xi} = \mathcal{O}\left(\frac{eM_{\mathrm{QCD}}}{(4\pi)^2}\right)\mathrm{Im}\,\Xi_{1,8}(M_{\mathrm{QCD}}). \tag{13}$$

The estimate for $d_n^{d_q}$ is in agreement with QCD sum-rule results [32, 33], while the estimate of d_n^{Ξ} agrees with the results of Refs. [34–36]. Combining Table II, Eq. (12) and (13) with the upper limit on the nEDM, $|d_n| \leq 2.9 \times 10^{-26}\, e\,\mathrm{cm}$ [37], and eEDM, $|d_e^{\mathrm{exp}}| \leq 8.7 \times 10^{-29}\, e\,\mathrm{cm}$ [38], we finally obtain the bounds shown in Table III.

V. CONCLUSION

Radiative β decay offers the possibility to study a *spin-independent* T-violating triple-correlation coefficient K. We have considered T-violating BSM physics arising above the

electroweak scale that contributes to this correlation. The dimension-6 operators that contribute to K also contribute to the spin-dependent EDMs. The EDM limits therefore stringently constrain these operators. In fact, comparing the EDM bounds in Table II to Eq. (9), we find that improving the EDM bounds would require a measurement of the neutron asymmetry better than 10^{-16}. This accuracy cannot be reached in present experiments. In conclusion, the T-odd correlation K is not "EDM-safe" when considering CP-violating dimension-6 operators above the electroweak scale.

ACKNOWLEDGMENTS

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